AXINO AS A STERILE NEUTRINO

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We present a supersymmetric axion model in which the fermionic superpartner of axion, i.e. the axino, corresponds to a sterile neutrino which would accommodate the LSND data with atmospheric and solar neutrino oscillations.

Current data from the atmospheric and solar neutrino experiments are beautifully explained by oscillations among three active neutrino species [1]. Another data in favor of neutrino oscillation has been obtained in the LSND experiment [2]. Reconciliation of these experimental results requires three distinct mass-squared differences, implying the existence of a sterile neutrino ν_s. In the four-neutrino oscillation framework, there are two possible scenarios [3,4]: the $2+2$ scheme in which two pairs of close mass eigenstates are separated by the LSND mass gap ~ 1 eV and the $3+1$ scheme in which one mass is isolated from the other three by the LSND mass gap. It has been claimed that the LSND results can be compatible with various short-baseline experiments only in the context of the 2+2 scheme. However, according to the new LSND results [2], the allowed parameter regions are shifted to smaller mixing angle, thereby allowing the $3+1$ scheme to be phenomenologically viable [3,4]. Although it can be realized in a rather limited parameter space, the 3+1 scheme is attractive since the fourth (sterile) neutrino can be added without changing the most favorable picture that the atmospheric and solar neutrino data are explained by the predominant $\nu_\mu \to \nu_\tau$ and $\nu_e \to \nu_\mu, \nu_\tau$ oscillations, respectively. In particular, the $3+1$ scheme with the heaviest ν_s would be an interesting explanation of all existing neutrino data. In this talk, we present a supersymmetric axion model with gauge-mediated supersymmetry (SUSY) breaking in which the fermionic superpartner of axion, i.e. the axino, corresponds to a sterile neutrino realizing the 3+1 scheme of four neutrino oscillation [5]. In this model, axino can be as light as 1 eV, and a proper axino-neutrino mixing is induced by R-parity violating couplings which appear as a consequence of spontaneous $U(1)_{PQ}$ breaking. It turns out that only the large angle MSW solution to the solar neutrino problem is allowed in this model.

The model under consideration contains three sectors: the observable sector, the SUSY-breaking sector, and the PQ sector. The observable sector contains the usual quarks, leptons, and two Higgs superfields, *i.e.* the superfields of the minimal supersymmetric standard model (MSSM). The SUSY-breaking sector contains a gauge-singlet Goldstino superfield X and the gauge-charged messenger superfields Y, Y^c as in the conventional gauge-mediated SUSY-breaking models [6]. Finally the PQ sector contains gauge-singlet superfields

S_k ($k = 1, 2, 3$) which break $U(1)_{PQ}$ by their vacuum expectation values (VEV), as well as gauge-charged superfields T, T^c which have the Yukawa coupling with some of S_k.

The Kähler potential of the model can always be written as

$$K = \sum_I \Phi_I^\dagger \Phi_I + ..., \tag{1}$$

where Φ_I denote generic chiral superfields of the model and the ellipsis stands for (irrelevant) higher dimensional operators which are suppressed by some powers of $1/M_*$ where M_* corresponds to the cutoff scale of the model which is presumed to be of order $M_{GUT} \sim 10^{16}$ GeV. The superpotential of the model is given by

$$W = hS_3(S_1 S_2 - f_{PQ}^2) + \kappa S_1 T T^c + \lambda X Y Y^c + W_{\rm MSSM} + W_{\rm SB} \tag{2}$$

where $W_{\rm MSSM}$ involves the MSSM fields, and $W_{\rm SB}$ describes SUSY breaking dynamics enforcing X develop a SUSY breaking VEV: $\langle \lambda X \rangle = M_X + \theta^2 F_X$. This VEV generates soft masses of the MSSM fields, $m_{\rm soft} \sim \alpha F_X / 2\pi M_X$, as in the conventional gauge-mediated SUSY breaking models [6]. One can easily arrange the symmetries of the model, e.g. $U(1)_{PQ}$ and an additional discrete symmetry, to make that $W_{\rm MSSM}$ is given by

$$\begin{aligned} W_{\rm MSSM} &= y_{ij}^{(E)} H_1 L_i E_j^c + y_{ij}^{(D)} H_1 Q_i D_j^c + y_{ij}^{(U)} H_2 Q_i U_j^c + \frac{y_0}{M_*} S_1^2 H_1 H_2 \\ &+ \frac{y_i'}{M_*^2} S_1^3 L_i H_2 + \frac{\gamma_{ijk}}{M_*} S_1 L_i L_j E_k^c + \frac{\gamma_{ijk}'}{M_*} S_1 L_i Q_j D_k^c + ..., \end{aligned} \tag{3}$$

where the Higgs, quark and lepton superfields are in obvious notations and the ellipsis stands for (irrelevant) higher dimensional operators.

To discuss the effective action at scales below f_{PQ}, let us define the axion superfield as

$$A = (\phi + ia) + \theta \tilde{a} + \theta^2 F_A,$$

where a, ϕ and $\tilde{a}$ are the axion, saxion and axino, respectively. It is then convenient to parameterize S_1 and S_2 as $S_1 = S e^{A/f_{PQ}}$, $S_2 = S e^{-A/f_{PQ}}$. As will be discussed later, the VEV of $e^{\phi/f_{PQ}} = \sqrt{S_1/S_2}$ can be determined to be of order unity by SUSY-breaking effects. We then have $\langle S_1 \rangle \approx \langle S_2 \rangle \approx f_{PQ}$, and then f_{PQ} corresponds to the axion decay constant which would determine most of the low energy dynamics of axion. After integrating out the SUSY-breaking sector as well as the heavy fields in the PQ sector, the low energy effective action includes the following Kähler potential and superpotential of the axion superfield A,

$$\begin{aligned} K_{\rm eff} &= f_{PQ}^2 \{ e^{(A+A^\dagger)/f_{PQ}} + e^{-(A+A^\dagger)/f_{PQ}} \} + \Delta K_{\rm eff}, \\ W_{\rm eff} &= \mu_0 e^{2A/f_{PQ}} H_1 H_2 + \mu_i' e^{3A/f_{PQ}} L_i H_2 \\ &\quad + e^{A/f_{PQ}} (\lambda_{ijk} L_i L_j E_k^c + \lambda_{ijk}' L_i Q_j D_k^c), \end{aligned} \tag{4}$$

where $\Delta K_{\rm eff}$ is *A-dependent* loop corrections involving the SUSY-breaking effects and

$$\begin{aligned}
&\mu_0 = y_0 f_{PQ}^2/M_*, \quad \mu_i' = y_i' f_{PQ}^3/M_*^2, \\
&\lambda_{ijk} = \gamma_{ijk} f_{PQ}/M_*, \quad \lambda'_{ijk} = \gamma'_{ijk} f_{PQ}/M_*.
\end{aligned} \tag{5}$$

The best lower bound on f_{PQ} is from astrophysical arguments implying $f_{PQ} \gtrsim 10^9$ GeV [7]. To accommodate the LSND data, we need the axino-neutrino mixing mass of order 0.1 eV. It turns out that this value is difficult to be obtained for $f_{PQ} > 10^{10}$ GeV. We thus assume $f_{PQ} = 10^9 - 10^{10}$ GeV with $M_* = M_{GUT}$ for which μ_0 takes an weak scale value (with appropriate value of y_0) [8] and the R-parity violating couplings $\lambda_{ijk}, \lambda'_{ijk}$ are appropriately suppressed. It is also easy to make the coefficient λ''_{ijk} of B and R-parity violating operators $U_i^c D_j^c D_k^c$ to be suppressed enough to avoid a too rapid proton decay into light gravitino and/or axino [9].

Low energy properties of the axion superfield crucially depends on how the saxion component is stabilized. One dominant contribution to the saxion effective potential comes from $\Delta K_{\rm eff}$ which is induced mainly by the threshold effects of T, T^c having the A-dependent mass $M_T = \kappa f_{PQ} e^{A/f_{PQ}}$. If $M_T \lesssim M_X$, one finds [10]

$$\Delta K_{\rm eff} \approx \frac{N_T}{16\pi^2} \frac{M_T M_T^\dagger}{\mathcal{Z}_T \mathcal{Z}_{T^c}} \ln\left(\frac{\Lambda^2 \mathcal{Z}_T \mathcal{Z}_{T^c}}{M_T M_T^\dagger}\right), \tag{6}$$

where N_T is the number of chiral superfields in T, $\mathcal{Z}_T$ is the Kähler metric of T, and Λ is a cutoff scale which is of order M_X. With $\mathcal{Z}_T|_{\theta^2\bar\theta^2} \approx -m_{\rm soft}^2$, $\Delta K_{\rm eff}$ of (6) gives a negative-definite saxion potential

$$V_\phi^{(1)} \approx -\frac{N_T}{16\pi^2} m_{\rm soft}^2 |\kappa f_{PQ}|^2 e^{2\phi/f_{PQ}}. \tag{7}$$

There is another (positive-definite) potential from the A-dependent μ-parameter:

$$V_\phi^{(2)} \approx e^{4\phi/f_{PQ}} |\mu_0|^2 (|H_1|^2 + |H_2|^2). \tag{8}$$

With $V_\phi^{(1)} + V_\phi^{(2)}$, the saxion can be stabilized at $\langle e^{\phi/f_{PQ}} \rangle \approx 1$ when κ is of order 10^{-6}. Once ϕ is stabilized at $\langle e^{\phi/f_{PQ}} \rangle \approx 1$, the resulting saxion and axino masses are given by

$$\begin{aligned}
m_\phi^2 &\approx (10 - 10^2\,{\rm keV})^2 + \Delta m_\phi^2, \\
m_{\tilde a} &\approx (10^{-4} - 10^{-2}\,{\rm eV}) + \Delta m_{\tilde a},
\end{aligned} \tag{9}$$

where the numbers in the brackets represent the gauge-mediated contributions for $f_{PQ} = 10^9 - 10^{10}$ GeV, $\mu_0 \approx 300$ GeV and $\langle e^{\phi/f_{PQ}} \rangle \approx 1$, and Δm_ϕ and $\Delta m_{\tilde a}$ are the supergravity-mediated contributions which are of order the gravitino mass $m_{3/2}$ as will be discussed in the subsequent paragraph.

The supergravity-mediated contributions to the saxion and axino masses can be quite model-dependent, in particular depends on the couplings of light moduli in the underlying supergravity model. However they are still generically of order the gravitino mass $m_{3/2}$ [11]. One model-independent supergravity-mediated contribution is from the auxiliary component u of the off-shell supergravity multiplet. In the Weyl-compensator formulation, u corresponds to the F-component of the Weyl compensator superfield:

$$\Phi = 1 + \theta^2 F_\Phi, \tag{10}$$

where the scalar component of Φ is normalized to be unity and the F-component is given by [12]

$$F_\Phi = e^{K/6}\left(m_{3/2} + \frac{F_I}{3}\frac{\partial K}{\partial \Phi_I}\right) \tag{11}$$

where $F_I = -e^{-K/2}K^{IJ}\partial(e^K W^\dagger)/\partial\Phi_J^\dagger$ denotes the F-component of Φ_I for the inverse Kähler metric K^{IJ} which is determined by the Kähler potential K of the underlying supergravity model. Note that Φ is defined as a dimensionless superfield, so F_Φ has mass-dimension one. The above expression shows that F_Φ is generically of order $m_{3/2}$. However it can be significantly smaller than $m_{3/2}$ in some specific models. For instance, in no-scale model with $K = -3\ln(T + T^\dagger - \Phi_i\Phi_i^\dagger)$ and $\partial W/\partial T = 0$, one easily finds $F_\Phi = 0$.

The Weyl-compensator contribution to the saxion and axino masses can be easily read off from the super-Weyl invariant supergravity action on superspace [12]:

$$-3\int d^4\theta\, \Phi\Phi^\dagger e^{-K/3} + [\int d^2\theta\, \Phi^3 W + \text{h.c.}] \tag{12}$$

This gives the following couplings of Φ to the axion superfield:

$$\int d^4\theta\, \Phi\Phi^\dagger K_{\text{eff}} = \int d^4\theta\, \Phi\Phi^\dagger (A + A^\dagger)^2 + ..., \tag{13}$$

where K_{eff} is the effective Kähler potential in (4). It is then straightforward to see that the Weyl compensator contributions to the saxion and axino masses are

$$\begin{aligned} \Delta m_\phi^2 &= 2|F_\Phi|^2 = \mathcal{O}(m_{3/2}^2), \\ \Delta m_{\tilde{a}} &= F_\Phi = \mathcal{O}(m_{3/2}). \end{aligned} \tag{14}$$

In gauge-mediated SUSY breaking models [6], the precise value of $m_{3/2}$ depends on the details of SUSY breaking sector. However most of models give $m_{3/2} \gtrsim 1$ eV, implying that $m_{\tilde{a}}$ of Eq. (9) is *dominated by the supergravity contribution* $\Delta m_{\tilde{a}}$. In this paper, we assume that $\Delta m_{\tilde{a}} \sim 1$ eV, so

$$m_{\tilde{a}} \approx 1 \text{ eV} \tag{15}$$

which would allow the axino to be a sterile neutrino for the LSND data. We note again that although it is generically of order $m_{3/2}$, $\Delta m_{\tilde{a}}$ can be significantly smaller than $m_{3/2}$ when the supergravity Kähler potential takes a particular form, e.g. the no-scale form [11].

Having defined our supersymmetric axion model, it is rather straightforward to compute the 4×4 axino-neutrino mass matrix:

$$\frac{1}{2} m_{\alpha\beta} \nu_\alpha \nu_\beta \tag{16}$$

where $\alpha, \beta = s, e, \mu, \tau$ and $\nu_s \equiv \tilde{a}$ with $m_{ss} = m_{\tilde{a}}$. The effective superpotential $W_{\rm eff}$ in (4) gives the following superpotential couplings

$$\int d^2\theta \left[\mu_0 (1 + \frac{2A}{f_{PQ}}) H_1 H_2 + \mu_i' (1 + \frac{3A}{f_{PQ}}) L_i H_2 \right]. \tag{17}$$

We will work in the field basis in which $\mu_i' L_i H_2$ ($i = e, \mu, \tau$) in $W_{\rm eff}$ are *rotated away* by an appropriate unitary rotation of H_1 and L_i. After this unitary rotation, the above superpotential couplings are changed to

$$\int d^2\theta \left[\mu_0 (1 + \frac{2A}{f_{PQ}}) H_1 H_2 + \frac{\mu_i' A}{f_{PQ}} L_i H_2 \right], \tag{18}$$

leading to the axino-neutrino mass mixing

$$m_{is} = \frac{\epsilon_i \mu_0 \langle H_2 \rangle}{f_{PQ}} \approx 0.1 \left(\frac{\epsilon_i}{10^{-5}} \right) \left(\frac{\mu_0}{600\,{\rm GeV}} \right) \left(\frac{10^{10}\,{\rm GeV}}{f_{PQ}} \right) {\rm eV}, \tag{19}$$

where $\epsilon_i = \mu_i' / \mu_0$.

The 3×3 mass matrix of active neutrinos is induced by R-parity violating couplings. At tree-level,

$$m_{ij} \approx \frac{g_a^2 \langle \tilde{\nu}_i^\dagger \rangle \langle \tilde{\nu}_j^\dagger \rangle}{M_a}, \tag{20}$$

where M_a denote the gaugino masses. The sneutrino VEV's $\langle \tilde{\nu}_i \rangle$ are determined by the bilinear R-parity violations in the SUSY-breaking scalar potential: $m^2_{L_i H_1} L_i H_1^\dagger + B_i' L_i H_2$. In our model, nonzero values of $m^2_{L_i H_1}$ and B_i' at the weak scale arise through renormalization group evolution (RGE), mainly by the coupling $\lambda'_{i33} y_b$ where y_b is the b-quark Yukawa coupling [13]. Moreover, $B H_1 H_2$ arises also through RGE which predicts a large $\tan\beta \approx 40 - 60$ [14]. We then find [13]

$$m_{ij} \approx 10^{-2} t^4 \left(\frac{\lambda'_{i33} y_b}{10^{-6}} \right) \left(\frac{\lambda'_{j33} y_b}{10^{-6}} \right) {\rm eV} \tag{21}$$

where $t = \ln(M_X / m_{\tilde{l}}) / \ln(10^3)$ for the slepton mass $m_{\tilde{l}}$. Here we have taken $m_{\tilde{l}} \approx 300$ GeV and $\mu_0 \approx 2 m_{\tilde{l}}$ which has been suggested to be the best parameter range for correct electroweak symmetry breaking [14].

Let us see how nicely all the neutrino masses and mixing parameters are fitted in our framework. The analysis of Ref. [4] leads to the four parameter regions, R1–R4 of Table I, accommodating the LSND with short baseline results. In our model, Eqs. (15) and (19) can easily produce the LSND mass eigenvalue $m_4 \approx m_{ss} = m_{\tilde{a}} \sim 1$ eV and also the LSND oscillation amplitude

$$A_{LSND} = 4U_{e4}^2 U_{\mu 4}^2 \approx 4 \left(\frac{m_{es}}{m_{ss}}\right)^2 \left(\frac{m_{\mu s}}{m_{ss}}\right)^2 \tag{22}$$

as the four mixing elements $U_{\alpha 4}$ of the 4×4 mixing matrix U are given by $U_{i4} \approx m_{is}/m_{ss} \approx 0.1$ $(i = e, \mu, \tau)$ and $U_{s4} \approx 1$. The masses and mixing of three active neutrinos can be easily analyzed by constructing the *effective* 3×3 mass matrix given by

$$m_{ij}^{\rm eff} = m_{ij} - \frac{m_{is} m_{js}}{m_{ss}}. \tag{23}$$

Upon ignoring the small loop corrections, this mass matrix has rank two, and can be written as

$$m_{ij}^{\rm eff} = m_{\rm x} \hat{\rm x}_i \hat{\rm x}_j + m_{\rm y} \hat{\rm y}_i \hat{\rm y}_j \tag{24}$$

where $\hat{\rm x}_i$ and $\hat{\rm y}_j$ are the unit vectors in the direction of m_{is} and $\langle \tilde{\nu}_j \rangle$, respectively. Remarkably, the mass scale $m_{\rm x} \approx (m_{is}/m_{ss})^2 m_{ss} \sim 10^{-2}$ eV gives the right range of the atmospheric neutrino mass. Eq. (21) shows that $m_{\rm y}$ is also in the range of 10^{-2} eV, so $m^{\rm eff}$ would be able to provide the right solar neutrino mass *unless* $\Delta m_{sol}^2 \ll 10^{-4}$ eV^2. Note from Eq. (5) that the typical size of $\epsilon_i, \lambda_{ijk}, \lambda'_{ijk}$ is around 10^{-6} for $f_{PQ} \approx 10^{10}$ GeV and $M_* \approx 10^{16}$ GeV.

The effective mass matrix $m_{ij}^{\rm eff}$ gives the mass eigenvalues

$$m_{2,3} = \frac{1}{2}\left(m_{\rm x} + m_{\rm y} \pm \sqrt{(m_{\rm x} + m_{\rm y}\cos^2 2\xi)^2 + m_{\rm y}^2 \sin^2 2\xi}\right) \tag{25}$$

and also the 3×3 mixing matrix of active neutrinos

$$U_{3\times 3} = (\hat{\rm z}^T, \hat{\rm w}^T c_\theta - \hat{\rm x}^T s_\theta, \hat{\rm w}^T s_\theta + \hat{\rm x}^T c_\theta)\,. \tag{26}$$

where $\hat{\rm z} \equiv \hat{\rm x} \times \hat{\rm y}/|\hat{\rm x} \times \hat{\rm y}|$, $\hat{\rm w} \equiv \hat{\rm x} \times \hat{\rm z}/|\hat{\rm x} \times \hat{\rm z}|$, $c_\xi \equiv \cos\xi = \hat{\rm x} \cdot \hat{\rm y}$ and $\tan 2\theta = m_{\rm y} \sin 2\xi/(m_{\rm x} + m_{\rm y}\cos 2\xi)$. The Super-Kamiokande data [15] combined with the CHOOZ result [16] imply that $U_{\mu 3}^2 \approx U_{\tau 3}^2 \approx 1/2$ and $U_{e3}^2 \ll 1$. The solutions to the solar neutrino problem can have either a large mixing angle (LA): $U_{e1}^2 \approx U_{e2}^2 \approx 1/2$, or a small mixing angle (SA): $U_{e1}^2 \approx 1$. This specify the first column $\hat{\rm z}^T$ of $U|_{3\times 3}$ as

$$({\rm LA}):\ \hat{\rm z} \approx (1/\sqrt{2}, -1/2, 1/2), \quad ({\rm SA}):\ \hat{\rm z} \approx (1, 0, 0)$$

up to sign ambiguities. Since $\hat{\rm x} \cdot \hat{\rm z} = 0$, the pattern $\hat{\rm z} \approx (1, 0, 0)$ implies $\hat{\rm x}_e \approx 0$. This leads to a too small $U_{e4} \approx m_{es}/m_{ss} \lesssim 10^{-2}$, so the SA solution is *not allowed* within our model. Among various LA solutions to the solar neutrino problem, *only* the large-angle MSW solution with $\Delta m_{sol}^2 \sim 10^{-4}\,{\rm eV}^2$ can be

naturally fitted since $m_{\rm x} \approx m_{\rm y} \sim 10^{-2}$ eV in our scheme. It is remarkable that $f_{PQ} \approx 10^{10}$ GeV and $M_* \approx M_{GUT}$ lead to the right size of R-parity violation yielding the desired values of m_{is} and m_{ij} also for the atmospheric and solar neutrino masses.

To see the feasibility of our whole scheme, we scanned our parameter space which consists of $m_{ss}, m_{is}, \lambda'_{i33} y_b$ to reproduce the allowed LSND islands R1–R4 of Table I together with the atmospheric and solar neutrino parameters [5]. For R1 and R4, we could find some limited parameter spaces which produce the corresponding oscillation parameters, however they need a strong alignment between $\hat{\rm x}$ and $\hat{\rm y}$ and also a large cancellation between $m_{\rm x}$ and $m_{\rm y}$. On the other hand, R2 and R3 do not require any severe fine tuning of parameters, so a sizable range of the parameter space can fit the whole oscillation data.

To conclude, we have shown that the 3+1 scheme of four-neutrino oscillation can be nicely obtained in supersymmetric axion model with gauge-mediated supersymmetry breaking. In this model, axino plays the role of sterile neutrino by having a mass ~ 1 eV and also a proper axino-neutrino mixing induced by R-parity violating couplings. One interesting feature of the model is that only the large angle MSW solution to the solar neutrino problem is allowed in this model.

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	$\lvert\Delta m^2_{41}\rvert(\mathrm{eV}^2)$	$\lvert U_{e4}\rvert$	$\lvert U_{\mu 4}\rvert$
R1	0.21-0.28	0.077-0.1	0.56-0.74
R2	0.88-1.1	0.11-0.13	0.15-0.2
R3	1.5-2.1	0.11-0.16	0.09-0.14
R4	5.5-7.3	0.13-0.16	0.12-0.16

Table 1. Allowed regions for the LSND oscillation.

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